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A complete calculation for direct detection of Wino dark matter

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Abstract

In the anomaly-mediated supersymmetry (SUSY) breaking scenario, neutral gaugino of $SU(2)_L$ multiplet, Wino, can be the lightest SUSY particle and become a candidate for dark matter. We calculated scattering cross section of Wino dark matter with nucleon, which is responsible for direct detection of the dark matter, on the assumption that the SUSY particles and the heavier Higgs bosons have masses of the order of the gravitino mass in the SUSY standard model. In such a case, the Wino-nucleon coupling is generated by loop processes. We have included two-loop contribution to Wino-gluon interaction in the calculation, since it is one of the leading contributions to the Wino-nucleon coupling. It was found that the spin-independent scattering cross section with proton is $10^{-(46-48)}$ cm^2. While it is almost independent of the Wino mass, the result is quite sensitive to the Higgs boson mass due to the accidental cancellation.

1 Introduction

The anomaly mediation [1] is the most economical mechanism to generate supersymmetry (SUSY) breaking terms in the supersymmetric standard model (SUSY SM). In the breaking mechanism, only dynamical SUSY-breaking sector is required, and no other extra fields are needed. On the assumption of the generic form of Kähler potential, all the scalar bosons except the lightest Higgs boson acquire masses, which are of the order of the gravitino mass. The gaugino masses, on the other hand, are generated by the quantum effects, and then they are suppressed by the one-loop factor compared with the gravitino mass. This is a concrete realization of the split SUSY scenario [2], in which the squarks and sleptons are $O(10^{(1-2)})$ TeV while the gaugino masses are less than $O(1)$ TeV. Such a mass spectrum is favored from phenomenological viewpoints of the SUSY flavor and CP problems [3] and the lightest Higgs mass bound [4]. Since it is safe from the cosmological gravitino over-production problem [5], it is also consistent with the thermal leptogenesis [6].

In the anomaly mediation the neutral component of $SU(2)_L$ gauginos, called as Winos, becomes the lightest in gaugino sector. This is because the gaugino masses are proportional to the beta functions of the gauge coupling. Higgsino, on the other hand, can be as heavy as gravitino, depending on the Kähler potential. Therefore, the neutral Wino can be the lightest SUSY particle (LSP) in the anomaly mediation scenario, and becomes a viable candidate for dark matter in the universe.

The thermal relic abundance of the Wino LSP in the universe is consistent with the WMAP observation when the Wino mass is from 2.7 TeV to 3.0 TeV [7]. The lighter Wino predicts too small thermal relic density; however, it is known that decay of gravitino or other quasi-stable particles may produce wino non-thermally so that the relic abundance is consistent with the observation [8, 9]. The successful Big-Bang Nucleosynthesis (BBN) also gives bounds on the annihilation cross section of the dark matter, while the large dark matter annihilation in the BBN era may give a solution to the lithium problem [10, 11]. In the anomaly mediation, the Wino mass around (150-300) GeV may be compatible with the lithium problem when the Wino LSP is the dominant component of the dark matter.

The direct detection of dark matter is now performed in several experiments with high sensitivities, and its theoretical sides are also extensively studied. The tree-level contribution to the Wino LSP-nucleon ($\tilde{\chi}^0$-N) scattering cross section, which is responsible for direct detection of the dark matter, is evaluated at Ref. [12]. However, in the case that the SUSY particles and the heavier Higgs bosons have masses of the order of the gravitino mass except the gauginos in the SUSY SM, the tree-level interactions of the Wino LSP with quarks are suppressed by the gravitino mass. Thus, the Wino LSP-nucleon scattering process is dominated by the weak gauge boson loop diagrams. However, despite the loop factor, it was pointed out that the loop contribution is not suppressed by the Wino mass even if it is heavier than the weak scale [13].

In this letter, we reevaluate the Wino LSP-nucleon scattering cross section. The one-loop contribution to the process is evaluated by Refs. [13, 14, 15]; however, their results are not consistent with each other. In addition, while the interaction of Wino and gluon is

generated by two-loop diagrams, it has to be included for the complete evaluation of the spin-independent Wino LSP-nucleon interaction. We take into account all the relevant diagrams up to two-loop and derive effective operators, which act as leading contribution in the scattering process.

2 Effective interaction for Wino LSP-nucleon scattering

First, we summarize the effective interactions of the Wino LSP with light quarks ($q = u, d, s$) and gluon, which are relevant to the Wino LSP-nucleon scattering. They are given as follows,

$$\mathcal{L}^{\rm eff} = \sum_{q=u,d,s} \mathcal{L}_q^{\rm eff} + \mathcal{L}_g^{\rm eff} , \tag{1}$$

where

$$\begin{aligned} \mathcal{L}_q^{\rm eff} &= d_q\, \overline{\tilde{\chi}^0}\gamma^\mu\gamma_5\tilde{\chi}^0\ \bar{q}\gamma_\mu\gamma_5 q + f_q m_q\ \overline{\tilde{\chi}^0}\tilde{\chi}^0\ \bar{q}q + f_q'\ \overline{\tilde{\chi}^0}\tilde{\chi}^0\ \bar{q}i\ \partial\!\!\!/ q \\ &+ \frac{g_q^{(1)}}{m_{\tilde{\chi}^0}}\ \overline{\tilde{\chi}^0}i\partial^\mu\gamma^\nu\tilde{\chi}^0\ \mathcal{O}^q_{\mu\nu} + \frac{g_q^{(2)}}{m^2_{\tilde{\chi}^0}}\ \overline{\tilde{\chi}^0}(i\partial^\mu)(i\partial^\nu)\tilde{\chi}^0\ \mathcal{O}^q_{\mu\nu} , \end{aligned} \tag{2}$$

$$\begin{aligned} \mathcal{L}_g^{\rm eff} &= f_G\ \overline{\tilde{\chi}^0}\tilde{\chi}^0 G^a_{\mu\nu}G^{a\mu\nu} \\ &+ \frac{g_G^{(1)}}{m_{\tilde{\chi}^0}}\ \overline{\tilde{\chi}^0}i\partial^\mu\gamma^\nu\tilde{\chi}^0\ \mathcal{O}^g_{\mu\nu} + \frac{g_G^{(2)}}{m^2_{\tilde{\chi}^0}}\ \overline{\tilde{\chi}^0}(i\partial^\mu)(i\partial^\nu)\tilde{\chi}^0\ \mathcal{O}^g_{\mu\nu} . \end{aligned} \tag{3}$$

Here, $m_{\tilde{\chi}^0}$ and m_q is mass of Wino and quark, respectively. The first term of $\mathcal{L}_q^{\rm eff}$ contributes to the spin-dependent $\tilde{\chi}^0$-N interaction, while the other terms in $\mathcal{L}_q^{\rm eff}$ and $\mathcal{L}_g^{\rm eff}$ generate spin-independent ones. The fourth and fifth terms in $\mathcal{L}_q^{\rm eff}$ and the second and third terms in $\mathcal{L}_g^{\rm eff}$ depend on the twist-2 operators (traceless parts of the energy momentum tensor) for quarks and gluon,

$$\begin{aligned} \mathcal{O}^q_{\mu\nu} &\equiv \frac{1}{2}\bar{q}i\left(\partial_\mu\gamma_\nu + \partial_\nu\gamma_\mu - \frac{1}{2}g_{\mu\nu}\ \partial\!\!\!/\right) q , \\ \mathcal{O}^g_{\mu\nu} &\equiv \left(G^{a\rho}_\mu G^a_{\rho\nu} + \frac{1}{4}g_{\mu\nu}G^a_{\alpha\beta}G^{a\alpha\beta}\right) . \end{aligned} \tag{4}$$

The scattering cross section of the Wino LSP with target nuclei is expressed compactly by using the coefficients given in $\mathcal{L}_q^{\rm eff}$ and $\mathcal{L}_g^{\rm eff}$ as follows [16],

$$\sigma = \frac{4}{\pi}\left(\frac{m_{\tilde{\chi}^0}m_T}{m_{\tilde{\chi}^0}+m_T}\right)^2\left[(n_pf_p+n_nf_n)^2 + 4\frac{J+1}{J}\left(a_p\langle S_p\rangle + a_n\langle S_n\rangle\right)^2\right] , \tag{5}$$

where m_T is the mass of target nucleus. The first term in the bracket comes from the spin-independent interactions while the second one is generated by the spin-dependent one. In

the spin-independent interaction term, n_p and n_n are proton and neutron numbers in the target nucleus, respectively, and the spin-independent coupling of the Wino with nucleon, f_N $(N = p, n)$, is given as

$$\begin{aligned} f_N/m_N &= \sum_{q=u,d,s} \left((f_q + f'_q) f_{Tq} + \frac{3}{4}(q(2) + \bar{q}(2))(g_q^{(1)} + g_q^{(2)}) \right) \\ &- \frac{8\pi}{9\alpha_s} f_{TG} f_G + \frac{3}{4} G(2) \left(g_G^{(1)} + g_G^{(2)} \right) \ . \end{aligned} \tag{6}$$

The matrix elements of nucleon are expressed by using nucleon mass m_N $(N = p, n)$ as[1]

$$\begin{aligned} f_{Tq} &\equiv \langle N | m_q \bar{q} q | N \rangle / m_N \ , \\ f_{TG} &\equiv 1 - \sum_{u,d,s} f_{Tq} \ , \\ \langle N(p) | \mathcal{O}^q_{\mu\nu} | N(p) \rangle &= \frac{1}{m_N} (p_\mu p_\nu - \frac{1}{4} m_N^2 g_{\mu\nu}) \ (q(2) + \bar{q}(2)) \ , \\ \langle N(p) | \mathcal{O}^g_{\mu\nu} | N(p) \rangle &= \frac{1}{m_N} (p_\mu p_\nu - \frac{1}{4} m_N^2 g_{\mu\nu}) \ G(2) \ . \end{aligned} \tag{7}$$

Here, $q(2)$, $\bar{q}(2)$ and $G(2)$ are the second moments of the quark, anti-quark and gluon distribution functions, which are expressed as

$$\begin{aligned} q(2) + \bar{q}(2) &= \int_0^1 dx \ x \ [q(x) + \bar{q}(x)] \ , \\ G(2) &= \int_0^1 dx \ x \ g(x) \ . \end{aligned} \tag{8}$$

They are scale-dependent, and are mixed with each others once the QCD radiative corrections are included. We use the second moments for gluon and quark distribution functions at the scale of Z boson mass, which are derived by the CTEQ parton distribution [17], and include bottom and charm quark contributions. On the other hand, the constant a_N $(N = p, n)$, which is responsible for the spin-dependent contribution, is defined as

$$a_N = \sum_{q=u,d,s} d_q \Delta q_N \ , \tag{9}$$

$$2 s_\mu \Delta q_N \equiv \langle N | \bar{q} \gamma_\mu \gamma_5 q | N \rangle \ , \tag{10}$$

where s_μ is the nucleon's spin, while J and $\langle S_N \rangle = \langle A | S_N | A \rangle$ in Eq. (5) are total spin of nucleus A and the expectation values of the total spin of protons and neutrons in A, respectively.

[1] We use equations of motion for quarks for evaluation of the matrix elements of $\langle N | \bar{q} i \not{\partial} q | N \rangle$, though this term is not relevant to our calculation, which we will see in the next section.

For proton	
f_{Tu}	0.023
f_{Td}	0.034
f_{Ts}	0.025
For neutron	
f_{Tu}	0.019
f_{Td}	0.041
f_{Ts}	0.025

Spin fraction	
Δu	0.77
Δd	-0.49
Δs	-0.15

Second moment at $\mu = m_Z$			
$G(2)$	0.48		
$u(2)$	0.22	$\bar{u}(2)$	0.034
$d(2)$	0.11	$\bar{d}(2)$	0.036
$s(2)$	0.026	$\bar{s}(2)$	0.026
$c(2)$	0.019	$\bar{c}(2)$	0.019
$b(2)$	0.012	$\bar{b}(2)$	0.012

Table 1: Parameters for quark and gluon matrix elements used in this letter. f_{Ti} ($i = u, d, s$) is taken from the estimation in Refs. [18, 19]. The spin fractions for proton comes from Ref. [20]. Those for neutron are given by exchange of up and down quarks in the tables. The second moments for gluon and quark distribution functions are calculated at the scale $\mu = m_Z$ (m_Z is Z boson mass) using the CTEQ parton distribution [17].

Notice that the term proportional to f_G in the spin-independent coupling of $\tilde{\chi}^0$-N in Eq. (6) is divided by α_s. It comes from definition of the gluon contribution to nucleon mass, f_{TG}, and the trace anomaly of the energy momentum tensor as

$$m_N f_{TG} = -\frac{9\alpha_s}{8\pi}\langle N|G^a_{\mu\nu}G^{a\mu\nu}|N\rangle \tag{11}$$

at the leading order.[2] Thus, when evaluating the spin-independent $\tilde{\chi}^0$-N interaction, we need to include $O(\alpha_s)$ correction to f_G [21]. Other contributions in the Wino LSP and gluon interaction, which come from gluon twist-2 operators, are sub-leading as far as the coefficients are $O(\alpha_s)$. Thus, we neglect the contribution from gluon twist-2 operators in the following discussion.

Parameters for quark and gluon matrix elements used in this analysis are summarized in Table 1. Notice that the strange quark content of the nucleon f_{Ts} is much smaller than previous thought according to the recent lattice simulation [19]. This leads to significant suppression on the spin-independent cross section, then the interaction of Wino and gluon becomes relatively more important in the cross section.

3 Results

Now we evaluate the coefficients of effective interactions in Eqs. (2, 3), which are needed to calculate the scattering cross section.

The Wino LSP accompanies the charged Wino ($\tilde{\chi}^-$). The mass difference is dominated by one-loop contribution unless Higgsino and Wino masses are almost degenerate; we ignore it in this letter. The coupling of neutral and charged Winos to the standard model

[2] Here, we use three-flavor approximation.

sector is only through gauge interactions as

$$\mathcal{L}_{\rm int} = -\frac{e}{s_W}\left(\overline{\tilde{\chi}^0}\gamma^\mu\tilde{\chi}^- W^\dagger_\mu + h.c.\right) + e\frac{c_W}{s_W}\overline{\tilde{\chi}^-}\gamma^\mu\tilde{\chi}^- Z_\mu + e\overline{\tilde{\chi}^-}\gamma^\mu\tilde{\chi}^- A_\mu\ . \tag{12}$$

Here, e is the electric charge, $s_W = \sin\theta_W$ and $c_W = \cos\theta_W$ with θ_W being the Weinberg angle. As is described in Introduction, the effective interactions of the Wino LSP to light quarks are generated by the loop diagrams. The leading contribution comes from one-loop interaction, which is shown in Fig. 1. After calculating the diagrams, the coefficients in Eq. (2) are derived as follows,

$$f_q = \frac{\alpha_2^2}{4m_W m_{h^0}^2} g_{\rm H}(x)\ , \tag{13}$$

$$f'_q = 0\ ,$$

$$d_q = \frac{\alpha_2^2}{m_W^2} g_{\rm AV}(x)\ , \tag{14}$$

$$g_q^{(1)} = \frac{\alpha_2^2}{m_W^3} g_{T1}(x)\ , \tag{15}$$

$$g_q^{(2)} = \frac{\alpha_2^2}{m_W^3} g_{T2}(x)\ , \tag{16}$$

where m_{h^0} is the lightest Higgs boson (*i.e.*, SM Higgs boson) mass, $x = m_W^2/m_{\tilde{\chi}^0}^2$ and $\alpha_2 = \alpha/s_W^2$ (here m_W is the W boson mass and α is the fine-structure constant). The diagram (a) in Fig. 1, which is induced by the SM Higgs boson (h^0) exchange, contributes to f_q, while the diagram (b) generates the other terms in Eq. (2). With the light quark masses ignored, the mass functions in Eqs. (13-16) are given as

$$\begin{aligned}
g_{\rm H}(x) &= -\frac{2}{b}(2+2x-x^2)\tan^{-1}(\frac{2b}{\sqrt{x}}) + 2\sqrt{x}(2 - x\log(x))\ ,\\
g_{\rm AV}(x) &= \frac{1}{24b}\sqrt{x}(8-x-x^2)\tan^{-1}(\frac{2b}{\sqrt{x}}) - \frac{1}{24}x(2-(3+x)\log(x))\ ,\\
g_{\rm T1}(x) &= \frac{1}{3}b(2+x^2)\tan^{-1}(\frac{2b}{\sqrt{x}}) + \frac{1}{12}\sqrt{x}(1-2x-x(2-x)\log x)\ ,\\
g_{\rm T2}(x) &= \frac{1}{4b}x(2-4x+x^2)\tan^{-1}(\frac{2b}{\sqrt{x}}) - \frac{1}{4}\sqrt{x}(1-2x-x(2-x)\log(x))\ ,
\end{aligned} \tag{17}$$

with $b = \sqrt{1-x/4}$.[3] As discussed in Ref. [13], the spin-independent interaction of $\tilde{\chi}^0$-N are not suppressed even if the Wino LSP is much larger than the W boson mass. The mass functions $g_{\rm H}(x)$ and $g_{\rm T1}(x)$ become finite in a limit of $x \to 0$ while other two functions

[3] Here, g_{T1} is larger than $F_{T1}^{(0)}$ given in Eq. (42) in [13]. We corrected it in this calculation.

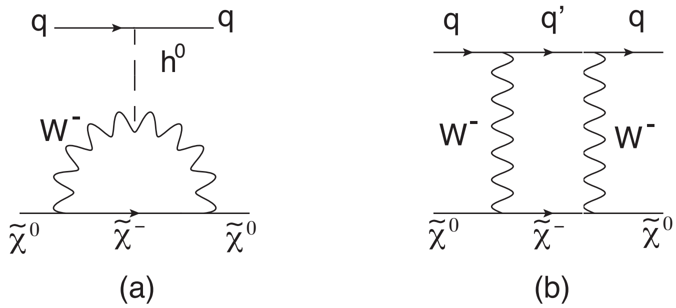

Figure 1: One-loop contributions to effective interactions of Wino LSP and light quarks.

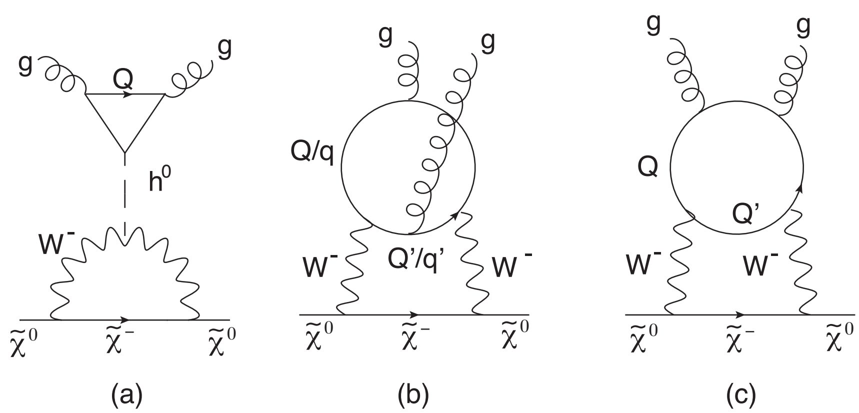

Figure 2: Two-loop contributions to interactions of Wino LSP and gluon. Here, Q and q represent heavy and light quarks, respectively.

are zero, as

$$
\begin{aligned}
g_{\rm H}(x) &\simeq -2\pi \ , \\
g_{\rm AV}(x) &\simeq \frac{\sqrt{x}}{6}\pi \ , \\
g_{\rm T1}(x) &\simeq \frac{\pi}{3} \ , \\
g_{\rm T2}(x) &\simeq -\frac{\sqrt{x}}{6} \ .
\end{aligned} \tag{18}
$$

Next, let us discuss the effective interactions of the Wino LSP and gluon. As we discussed in the previous section, the $O(\alpha_s)$ correction to f_G in Eq. (3) is relevant at the leading order though it is induced by two-loop order. Three types of diagrams in Fig. 2 contribute to f_G. The diagram (a) includes heavy quark loop ($Q = c, b, t$). The heavy quark content of the nucleon is related to the gluon condensate as [22]

$$
\langle N | m_Q \bar{Q} Q | N \rangle = -\frac{\alpha_s}{12\pi} \langle N | G^a_{\mu\nu} G^{a\mu\nu} | N \rangle \ . \tag{19}
$$

Thus, the diagram (a) can be evaluated from Eq. (13) by replacing light to heavy quarks and using Eq. (19). On the other hand, we need to calculate irreducible two-loop diagrams (b) and (c) explicitly. In the diagram (c), the momentum which dominates the quark loop integration is characterized by mass of quark which emits two gluons. Since we are constructing the effective theory under $O(1)$ GeV, the integration in the infrared regime under such energy scale should not be included. Thus, light quarks does not contribute in this diagram. On the other hand, the loop momentum of quark loop in the diagram (b) is dominated by the external momentum of the quark loop diagram (*i.e.*, W boson mass in this case); therefore, all quarks contributes in the loop. We express the $O(\alpha_s)$ contribution to f_G as follows,

$$f_G = -3 \times \frac{\alpha_s}{12\pi}\frac{\alpha_2^2}{4m_W m_{h^0}^2} g_{\rm H}(x) + \frac{\alpha_s}{4\pi}\frac{\alpha_2^2}{m_W^3} g_{\rm B3}(x,y) + 2\times\frac{\alpha_s}{4\pi}\frac{\alpha_2^2}{m_W^3} g_{\rm B1}(x) \ , \qquad (20)$$

where $y = m_t^2/m_{\tilde{\chi}^0}^2$ (m_t is top quark mass). The first term represents the contribution from the diagram (a). The second term comes from the diagrams (b) and (c) with the third-generation quark loop, while the the first- and second-generation quarks contribute to the third one. We ignore quark masses except for the top quark. Then, we found that the mass functions $g_{\rm B3}(x,y)$ and $g_{\rm B1}(x)$ are given by

$$\begin{aligned}
g_{\rm B3}(x,y) &= -\frac{x^{3/2}(2y-x)}{12(y-x)^2} - \frac{x^{3/2}y^3\log(y)}{24(y-x)^3} + \frac{x^{5/2}(3y^2-3xy+x^2)\log(x)}{24(y-x)^3} \\
&+ \frac{x^{3/2}\sqrt{y}(y^3-2y^2-14y+6x)\tan^{-1}(\frac{2b_t}{\sqrt{y}})}{24b_t(y-x)^3} \\
&- \frac{x\,(x^4-3yx^3-2x^3+3y^2x^2+6yx^2+4x^2-6y^2x-6yx-6y^2)\tan^{-1}(\frac{2b}{\sqrt{x}})}{24b(y-x)^3} \ , \\
g_{\rm B1}(x) &= -\frac{1}{24}\sqrt{x}(x\log(x)-2) + \frac{(x^2-2x+4)\tan^{-1}(\frac{2b}{\sqrt{x}})}{24b} \ ,
\end{aligned} \qquad (21)$$

where $b_t = \sqrt{1-y/4}$. Notice that the diagrams (b) and (c) also give finite contributions to the spin-independent $\tilde{\chi}^0$-N interaction in a limit of $m_{\tilde{\chi}^0} \to \infty$, *i.e.*, $x,\ y \ll 1$,

$$\begin{aligned}
g_{\rm B3}(x,y) &\simeq \frac{(3\sqrt{y}+2\sqrt{x})x}{24(\sqrt{x}+\sqrt{y})^3}\pi \ , \\
g_{\rm B1}(x) &\simeq \frac{\pi}{12} \ .
\end{aligned} \qquad (22)$$

Now we are at the position to present the scattering cross section. In Fig. 3, we show the spin-independent $\tilde{\chi}^0$-p scattering cross section as a function of $m_{\tilde{\chi}^0}$ (solid line). Here, we take m_{h^0} = 115, 130, 300 GeV, and 1 TeV from bottom to top. While the latter two values may not be realistic in the minimal SUSY SM, the next-minimal SUSY SM (NMSSM), for example, may predict larger Higgs boson mass. It was found that the spin-independent cross section is $O(10^{-(48-46)})$ cm^2, depending on the Higgs boson mass.

In order to understand the result, we also plot each contribution from the effective operators in f_p in Fig. 4. Solid line represents the Higgs exchange contribution (Fig. 1(a) and Fig. 2(a)), dashed line is for the twist-2 operator contribution (Fig. 1(b)), and dash-dot line is for that from irreducible two-loop diagrams in Fig. 2(b) and (c). As is seen, the contribution from quark twist-2 operators is dominant part. However, we also found that other two also give relatively large contribution by the opposite sign. Consequently, f_p is suppressed by the accidental cancellation, which leads to the smaller spin-independent cross section. When the Higgs boson mass is taken to be larger, the cross section becomes larger since the cancellation is milder.

In this letter, we have ignored the tree-level coupling of the Wino LSP and the lightest Higgs boson since it is suppressed by heavy Higgsino mass. When it dominates the spin-independent interaction, the spin-independent cross section is evaluated as

$$\sigma_p \simeq 9 \times 10^{-47}\mathrm{cm}^2 \times \left(\frac{m_{\tilde{H}}}{10\mathrm{TeV}}\right)^{-2} \left(\frac{m_{h^0}}{115\mathrm{GeV}}\right)^{-4} \sin^2 2\beta \ , \tag{23}$$

where $m_{\tilde{H}}$ is the Higgsino mass and β is the vacuum angle in the SUSY SM. Thus, the tree-level contribution may dominate the spin-independent cross section, depending on parameters in the SUSY SM, even if the the SUSY particle masses are of the order of the gravitino mass.

For completeness, in Fig. 3, we also show the spin-dependent cross section in dashed line. As expected from the behavior of the mass function $g_{\mathrm{AV}}(x)$, it was found that the cross section is suppressed by the Wino mass.

Finally, we comment difference between our result and the previous works. We found a few errors in the calculation in Ref. [13] as described in this text, though the qualitative behavior is not different from this work. On the other hand, compared with Refs. [14, 15], our result for the spin-independent cross section is smaller by $O(10^{-(2-3)})$. In Ref. [14] only the contribution to scalar couplings to quarks and gluon are evaluated. In Ref. [15] the relative sign of the quark twist-2 and the Higgs boson exchange contributions is opposite to ours. Thus, the cross section is not reduced by accidental cancellation in those works. Their loop functions are also different from ours. We could not understand origin of the differences.

4 Conclusion and discussion

In this letter, we calculated the Wino LSP-nucleon cross section in the anomaly-mediated SUSY breaking mechanism. We especially consider the scenario in which all SUSY particles except for gauginos are heavy to decouple in electroweak scale, and neutral Wino becomes the LSP. In such a scenario, although the Wino LSP does not interact with nucleon at tree-level, it does in loop diagrams. We have taken into account all the loop diagrams which act as leading contributions to the Wino LSP-nucleon scattering. As a result, the spin-independent cross section turns out to be $O(10^{-(48-46)})$ cm^2, depending on the Higgs boson mass. In the calculation, we found that Wino-gluon interaction at two-loop level contributes in opposite sign to the main Wino-quark interaction, which

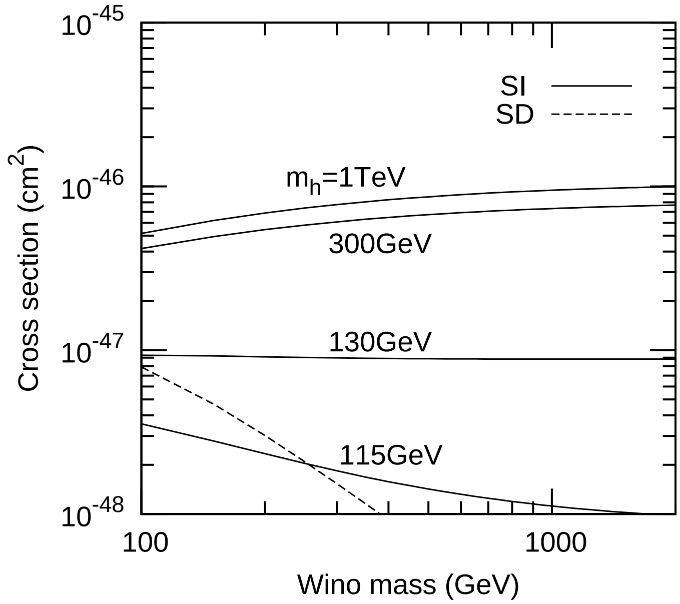

Figure 3: $\tilde{\chi}^0$-p scattering cross section as a function of $m_{\tilde{\chi}^0}$. Spin-independent (SI) cross section is given in solid line, taking m_{h^0} = 115, 130 GeV, 300 GeV, and 1 TeV from bottom to top. Here, we also plot spin-dependent (SD) one in dashed line.

leads to the suppression of the total Wino-nucleon coupling. Therefore, it is concluded that the direct detection of Wino dark matter is difficult in the present experiments in the scenario.

We comment on the cancellation that we observed in Wino-nucleon coupling. The cancellation is supposed to be accidental in the scenario that we analysed. Thus, if one consider the other scenarios in SUSY or other models, the two-loop may contribute as large as the lower order diagrams, which may cause enhancement of scattering cross section. Such analysis will be given elsewhere [23].

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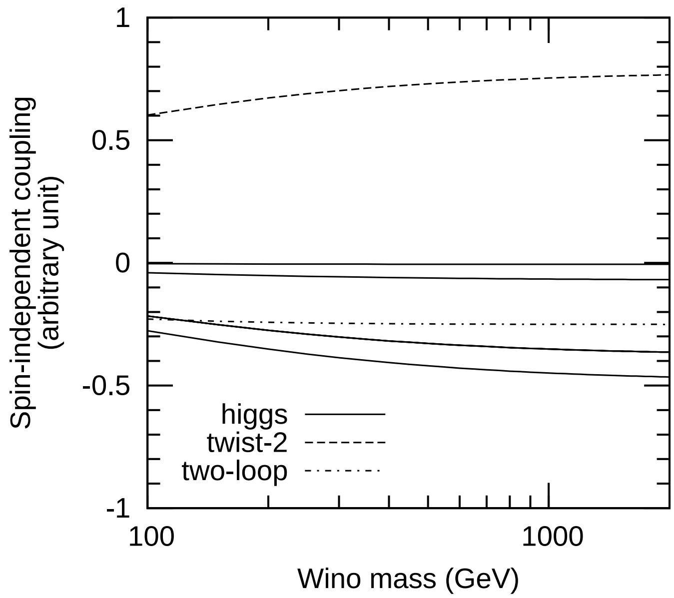

Figure 4: Each contribution in spin-independent coupling, f_p. Solid line represents the Higgs exchange contribution including heavy quark one, dashed line is for the twist-2 operator contribution, and dash-dot line is for that from two-loop diagrams in Fig. 1(b) and (c). The Higgs boson mass is $m_{h^0} = 115,\ 130$ GeV, 300 GeV, and 1 TeV from bottom to top. Here, unit is arbitrary.

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